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Optical-infrared flares and radio afterglows by Jovian planets inspiraling into their host stars

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ABSTRACT
When a planet inspirals into its host star, it releases gravitational energy which is converted into an expanding bubble of hot plasma. We study the radiation from the bubble and show that it includes prompt optical-infrared emission and a subsequent radio afterglow. The prompt emission from M31 and Large Magellanic Cloud is detectable by optical-near infrared transient surveys with a large field of view. The subsequent radio afterglows are detectable for 10^{3-4} years. The event rate depends on uncertain parameters in the formation and dynamics of giant planets. Future observation of the rate will constrain related theoretical models. If the event rate is high ($\gtrsim$ a few events per year), the circumstellar disk must typically be massive as suggested by recent numerical simulations.

Key words: radiation mechanisms: non-thermal — radiation mechanisms: thermal — planet-star interactions — infrared: planetary systems — radio continuum: planetary systems

1 INTRODUCTION

A substantial fraction of gaseous planetary-mass objects might be ingested by the central stars (Machida et al. 2010, 2011; Inutsuka 2012; Vorobyov & Basu 2010, 2015, and references therein). Realistic non-ideal magnetohydrodynamics simulations have shown that protoplanetary disks are initially massive enough to produce multiple Jupiter-mass planets via gravitational instability (Inutsuka et al. 2010; Machida et al. 2011; Tsukamoto et al. 2015). These massive planets may survive in the subsequent era, during which planets gravitationally interact or collide with each other to produce hot Jupiters and highly eccentric planets (e.g., Ida & Lin 2004; Chatterjee et al. 2008; Ford & Rasio 2008). A large fraction of the hot Jupiters which migrate to the vicinity of the central stars are either consumed (Sandquist et al. 1998) or tidally disrupted (Gu et al. 2003) by the host star. Stars without detected hot Jupiters might have already ingested them (Rice et al. 2008; Inutsuka 2009; Ogihara et al. 2013, 2014). Present-day hot Jupiters could secularly enlarge their eccentricity to reach their host stars by a process like the Kozai-Lidov mechanism (Kozai 1962; Lidov 1962). Several ways for detecting stellar ingestion of planets have been proposed (e.g., Sandquist et al. 1998, 2002; Cody & Sasselov 2005; Jackson et al. 2009; Teitler & Konigl 2014; Matsakos & Konigl 2015). If a star ingests planets on average $N_{\rm i}$ times during its life, the total event rate in the Galaxy is estimated to be ${\rm SFR}\times N_{\rm i}/\langle m\rangle \sim 5N_{\rm i}$ yr^{-1}, where SFR $\approx 1M_\odot$ yr^{-1} is the star formation rate of the Milky Way (Robitaille & Whitney 2010) and $\langle m\rangle \approx 0.2M_\odot$ is the average stellar mass (Kroupa 2001; Chabrier 2003).

In this paper, we calculate the radiation expected at the moment of Jovian planet ingestion and the subsequent afterglow phase. When a planet is engulfed by a host star, it releases gravitational energy which is converted into an expanding bubble of hot plasma (§ 2). The expanding bubble generates an optical-infrared flare (§ 3). Subsequently, the material interacts with the circumstellar matter (CSM), generating a shock that accelerates electrons to relativistic energies, which in turn produce synchrotron radiation (§ 4). There are several previous papers that predicted transient emissions from planet-star interaction. Metzger et al. (2012) focused on the case of quasi-circular orbit of the hot Jupiters around a main sequence star, which produces super-Eddington accretion. Bear et al. (2011) studied tidal disruption of planets by brown dwarfs. These authors argued that the accretion disk drives outflows, producing long-duration (> day) transients. On the other hand, we argue that a planet with a highly eccentric orbit in-spirals into the star, depositing thermal energy near the stellar surface, which produces the expanding plasma bubble.

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2 INTERACTION BETWEEN A STAR AND A JOVIAN PLANET

Most extrasolar planets reside in eccentric orbits (Marcy et al. 2005; Winn & Fabrycky 2015). Some planets are subject to rapid orbital change or migration by the interaction with a gaseous disk (Rice et al. 2008), planet-planet scattering (Rasio & Ford 1996; Weidenschilling & Marzari 1996), or the Kozai-Lidov mechanism (Ford et al. 2000; Fabrycky & Tremaine 2007; Li et al. 2015). Following these dynamical processes, the pericentre distance $R_{\rm p}$ finally becomes shorter than the tidal disruption radius of a planet: $R_{\rm T} = (M_*/m_{\rm pl})^{1/3} r_{\rm pl}$, where M_* is a mass of the star hosting the planet with a mass $m_{\rm pl}$ and a radius $r_{\rm pl}$. Then, the planet is tidally disrupted (Rees 1988; Phinney 1989; Evans & Kochanek 1989 for the context of a supermassive black hole system, Faber et al. 2005; Guillochon et al. 2011 for a star-planet system). Note that R_T is as large as the radius of the star R_* for Sun-like G-type stars and Jovian mass planets, whereas $R_T > R_*$ for K- and M-stars (e.g., Rappaport et al. 2013). If $R_{\rm p}$ is smaller than $R_{\rm T}$ for the Sun-like G-type star we consider in this paper, the approaching planet can impact directly the stellar surface. In other words, if a penetration factor $\beta = R_{\rm T}/R_{\rm p} = 1$, the planet is tidally disrupted by the host star. Otherwise, for $\beta > 1$ the planet would collide with the host star. Next, we will discuss these two cases separately in the following subsections.

2.1 Tidal disruption of a planet by a star: $\beta = 1$

After the tidal disruption of a planet on a parabolic orbit, the debris mass is distributed around zero specific energy (e.g., see Figure 3 of Evans & Kochanek 1989), characteristic of a parabolic orbit. If the planet originally approaches the star on a bound (eccentric or circular) orbit, however, the mass distribution of the disrupted planet would shift to negative specific energy. This gives the condition to cause tidal disruption of planets on eccentric orbits, $0 \leqslant e < e_{\rm crit}$, where $e_{\rm crit}$ is related to a penetration factor β (Hayasaki et al. 2013).

$$e_{\rm crit} = 1 - \left(\frac{2}{\beta}\right)\left(\frac{m_{\rm pl}}{M_*}\right)^{1/3} \quad . \tag{1}$$

On the other hand, $e_{\rm crit} < e \leqslant 1$ must be satisfied for the parabolic tidal disruption events (TDEs). In order for the planet to be tidally disrupted, the pericentre distance must be smaller than the tidal disruption radius. This gives a constraint on the orbital eccentricity:

$$e \geqslant 1 - \frac{R_{\rm T}}{a} \tag{2}$$

Combining Eqs. (1) with (2), we can obtain the alternative condition that $a_{\rm crit} < a < \infty$ for parabolic TDEs and $0 < a < a_{\rm crit}$ for the eccentric TDEs, where

$$a_{\rm crit} = \frac{\beta}{2}\left(\frac{M_*}{m_{\rm pl}}\right)^{1/3} R_{\rm T}. \tag{3}$$

For the Sun-like G-type stars with Jovian planets, the critical semi-major axis is estimated to be $a_{\rm crit} \approx 5\beta R_\odot$, so that parabolic TDEs only occur if $a \gtrsim a_{\rm crit}$ for $\beta = 1$. The eccentric TDEs are likely to happen in the case of the system composing of a star and a hot-jupiter on a quasi-circular orbit, because of $0 < a \lesssim a_{\rm crit}$[1].

Assuming the full conservation of angular momentum during debris circularization, the circularization radius is given by $r_{\rm c} = (R_{\rm T}/\beta)(1+e)$ for Keplerian rotation. If $\beta = 1$, an accretion disc forms around the star at $R_* \lesssim r_{\rm c} \lesssim 2R_*$ for $0 \lesssim e \lesssim 1$. Because of the super-Eddington accretion nature of mass fallback rate, the super-Eddignton outflow is likely to be caused (Strubbe & Quataert 2009). This outflow would produce the long-duration ($\gg$ day) optical transients occurred in the star-planet system (Metzger et al. 2012). On the other hand, if $\beta > 1$, the planet impacts the stellar surface and is subject to orbital friction within the stellar envelope. In this regime, it is not clear whether the tidal disruption occurs.

2.2 Direct collision between a star and a planet: $\beta > 1$

2.2.1 *Simple order-of-magnitude estimate*

Let us consider the situation that the planet on highly eccentric orbit hits the star. The ram pressure the planet gets from the star is given by $\rho_* v_{\rm p}^2$, where $\rho_* = (3/4\pi)M_*/R_*^3$ and $v_{\rm p} = \sqrt{GM_*/R_*}$ are the average stellar density and planetary velocity, respectively. Then, the drag force on the planet is given by,

$$F_{\rm drag} = \eta \sigma_{\rm p} \rho_* v_{\rm p}^2, \tag{4}$$

where $\eta = \mathcal{O}(1)$ and $\sigma_{\rm p}$ are the geometrical parameter and the cross section, respectively. For simplicity, we adopt $\sigma_{\rm p} = \pi r_{\rm pl}^2$. Our estimated drag force is in agreement with recent simulation results (Staff et al. 2016). If the planet travels a distance d inside the star, the total energy lost by the drag force $\Delta E_{\rm p}$ can be estimated as,

$$\Delta E_{\rm p} = F_{\rm drag}\cdot d = \frac{3}{4}\eta\frac{GM_*^2}{R_*}\left(\frac{r_{\rm pl}}{R_*}\right)^2\left(\frac{d}{R_*}\right) \quad . \tag{5}$$

The kinetic energy of the planet at the periastron is given by,

$$\begin{aligned} E_{\rm p} &= \frac{1}{2} m_{\rm pl} v_{\rm peri}^2 \sim \frac{\beta}{2}\frac{GM_* m_{\rm pl}}{R_*} \\ &= 3.6\times10^{45}\left(\frac{\beta}{2}\right)\left(\frac{\xi_*}{1}\right)\left(\frac{m_{\rm pl}}{m_{\rm J}}\right) {\rm erg}, \end{aligned} \tag{6}$$

where $v_{\rm peri} = \sqrt{GM_*/R_{\rm p}} \sim \sqrt{\beta GM_*/R_*}$ is the planetary velocity at the pericenter, $\xi_* = (M_*/M_\odot)/(R_*/R_\odot)$ and $m_{\rm J} = 1.898\times10^{30}$g is Jupiter's mass. Note that ξ_* is of order unity for main sequence stars with $M_* \lesssim 2M_\odot$ (e.g., Torres et al. 2010; Eker et al. 2015).

The ratio of $\Delta E_{\rm p}$ and $E_{\rm p}$ is given by

$$\frac{\Delta E_{\rm p}}{E_{\rm p}} = \frac{3}{2}\left(\frac{\eta}{\beta}\right)\left(\frac{M_*}{m_{\rm pl}}\right)\left(\frac{r_{\rm pl}}{R_*}\right)^2\left(\frac{d}{R_*}\right) \quad . \tag{7}$$

For a Sun-like G-type star with the Jovian planet and $\eta \sim 1$ and $\beta \sim 2$, $\Delta E_{\rm p}/E_{\rm p} \sim 10$ for $d = R_*$. This suggests that the planet spirals into the star and stops when $\Delta E_{\rm p} \approx E_{\rm p}$

[1] The tidal disruption of the planet given by Section 5 of Metzger et al. (2012) corresponds to the eccentric TDEs we have proposed here.

and $d \approx 0.1R_* \sim r_{\rm J}$, where $r_{\rm J} = 7.0 \times 10^9$ cm is the radius of Jupiter. Following the constraint on β for $d < r_{\rm peri}$ as $2 < \beta < \sqrt{15}$ from Eq. (7), we find that for moderate values of β the inspiraling planet could stop in the stellar external layer before causing the tidal disruption. The friction between the planet and the stellar gas deposits the thermal energy. If the supersonic planet motion makes a bow shock around the planet, dissipated energy also turns into the thermal energy. In these ways, independently of whether the tidal disruption occurs, the thermal energy,

$$E_{\rm th,i} \sim \frac{GM_* m_{\rm pl}}{2R_*} = 1.8 \times 10^{45} \xi_* \left(\frac{m_{\rm pl}}{m_{\rm J}}\right) \ {\rm erg} \quad , \tag{8}$$

is injected at the stellar external layer. The gas is optically thick, so that the radiative cooling is inefficient, making the hot, adiabatically expanding plasma bubble.

2.2.2 An example of star-planet collision: the case of Sun-like star and a Jupiter-like planet.

Below we consider the collision between a Sun-like star and a Jupiter-like giant planet as a typical example. We start by deriving simple power-law scaling relations for the solar interior. For a geometrically thin adiabatic gas with an adiabatic index γ_g, hydrostatic equilibrium gives the temperature as a function of the vertical depth z (distance from the solar surface) as $T_*(z) = \mu\mu_{\rm p}(\gamma_g - 1)gz/k_B\gamma_g$, where g, $\mu_{\rm p}$ and μ are the gravitational acceleration, the proton mass and the mean molecular weight, respectively. In terms of the polytropic index $n = (\gamma_g - 1)^{-1}$, we find that the mass density and pressure scale as $\rho_* \propto z^n$ and $p_* \propto z^{n+1}$. Comparing these scalings with standard solar models (Guenther et al. 1992), we find that the simple power-law expressions with $n = 2$ ($\gamma_g = 3/2$),

$$T_*(z) = 1 \times 10^6 \ {\rm K} \left(\frac{z}{10^{10}{\rm cm}}\right) \quad , \tag{9}$$

$$\rho_*(z) = 0.05 \ {\rm g \ cm^{-3}} \left(\frac{z}{10^{10}{\rm cm}}\right)^2 \quad , \tag{10}$$

$$p_*(z) = 5 \times 10^{12} \ {\rm dyn \ cm^{-2}} \left(\frac{z}{10^{10}{\rm cm}}\right)^3 \quad , \tag{11}$$

approximate the solar interior well for $z \lesssim 2 \times 10^{10}$cm. Note that main sequence stars with a mass larger than $1M_\odot$ are more concentrated than $n = 3$ polytropes (see, for example, Figure A1 of Freitag & Benz 2005).

The deceleration of the planet inside a Sun-like star is described by[2]

$$\frac{{\rm d}\ln v_{\rm p}}{{\rm d}z} = -\frac{\eta\rho_*(z)}{\rho_{pl} r_{pl}\cos\theta} \quad , \tag{12}$$

where $\rho_{pl} = 3m_{pl}/4\pi r_{pl}^3$ is the average mass density of the planet, η is a numerical factor of order of unity [see Eq. (4)], and θ is an angle between the surface normal and the velocity of the planet just before the collision, that is, $z = d\cos\theta$. With the initial velocity $v_{\rm p0}$ at $z = 0$, equations (10) and (12) give

$$v_{\rm p}(z) = v_{\rm p0} \exp\left[-\left(\frac{z}{z_s}\right)^3\right] \quad , \tag{13}$$

where the stopping depth z_s is given by

$$z_s = 3.8 \times 10^{10} \ {\rm cm} \left(\frac{\cos\theta}{\eta}\right)^{1/3} \left(\frac{m_{\rm pl}}{m_{\rm J}}\right)^{1/3} \left(\frac{r_{\rm pl}}{r_{\rm J}}\right)^{-2/3} . \tag{14}$$

Hence the planet keeps its velocity until $z \lesssim z_s$, and almost suddenly stops at a depth z_s. For a Sun-like density structure, the density more rapidly increases for $z \gtrsim 2 \times 10^{10}$ cm, so that the rapid deceleration starts before z_s. Note that in deriving Eq. (14), several effects were neglected for the purpose of a simple analytical calculation. The sound speed of the stellar gas is

$$\begin{aligned} c_{s*}(z) &= \left(\frac{\gamma_g p_*(z)}{\rho_*(z)}\right)^{1/2} \\ &\approx 1 \times 10^7 \ {\rm cm \ s^{-1}} \left(\frac{z}{10^{10}{\rm cm}}\right)^{5/2} \quad , \end{aligned} \tag{15}$$

so that for the planet with initial velocity $v_{\rm p0} = 6.2 \times 10^7$ cm s^{-1} (which is the escape velocity of the sun), the Mach number is

$$\mathcal{M} = \frac{v_p(z)}{c_{s*}(z)} \approx 5 \left(\frac{z}{10^{10}{\rm cm}}\right)^{-5/2} \exp\left[-\left(\frac{z}{z_s}\right)^3\right] \quad , \tag{16}$$

Therefore, the planet motion inside the star is mildly supersonic and a bow shock is formed. As a result, the planet mass decreases due to ablation. The ram pressure of the stellar gas also causes lateral expansion of the planet, which increases the cross sectional area $\sigma_{\rm p}$ [see Eq. (4)]. It is difficult to treat these effects analytically, but they effectively increase the value of η, so that z_s becomes smaller. Planets could also inflate as they approach the star due to tidal heating. Recent observations have shown that hot Jupiters have a larger radius than expected for given planet mass (e.g Baraffe et al. 2014). Taking into account all of these additional effects beyond our simple estimate in Eq. (14), the actual stopping depth may be a factor of a few smaller, say $z_s \approx 1 - 2 \times 10^{10}$ cm. This is of order $r_{\rm J} \approx 0.1R_\odot$, in agreement with the previous simple estimate given by Eq. (7) in § 2.2.1. Around this depth, the planet's kinetic energy is suddenly released, so that the thermal energy, $E_{\rm th,i}$, is deposited there. The hot bubble arises there. The bubble temperature is $T_i \sim 10^{6-7}$ K, based on Eq. (17), and is slightly larger than the temperature of the stellar gas at z_s, $T_*(z_s)$.

Following the collision, the expanding bubble rises and finally escapes the stellar surface. The expansion velocity is of order the sound speed of the bubble, $\sqrt{k_B T_i/\mu_p}$, which is of the order the sound speed of the stellar gas at the base, $c_{s*}(z_s)$. The expansion time of the bubble is $t_{\rm exp}(z) \sim z/c_{s*}(z_s)$. This timescale is comparable to the time that the stellar gas material fills the rarefied region behind the planet, $t_{\rm cl}(z) \sim r_{\rm pl}/c_{s*}(z)$. Hence, we obtain $t_{\rm exp}(z)/t_{\rm cl}(z) \sim (z/r_{\rm pl})[c_{s*}(z)/c_{s*}(z_s)] = (z_s/r_{\rm pl})(z/z_s)^{7/2}$. Since z_s is slightly larger than $r_{\rm pl}$, the rarefied region closes before the bubble expands. However, a strong pressure wave will travel to the surface of the star and lift material from there out of the gravitational potential well of the star. The actual impact of a collision can only be reliably calculated with a numerical hydrodynamics simulation, which we leave for future work.

[2] For a similar study of a Jupiter-comet collision, see Chevalier & Sarazin (1994).

3 PROMPT EMISSION FROM EXPANDING PLASMA BUBBLE

As seen in § 2.2, the collision between a star and a Jovian planet releases thermal energy $E_{\rm th,i}$ into a volume of radius $R_{\rm i}$ over a short time. We approximate $R_{\rm i} \sim r_{\rm pl}$ and the number density of the confined gas, $n_{\rm i} \sim 3m_{\rm pl}/4\pi\mu_{\rm p}R_{\rm i}^3$, where $\mu_{\rm p}$ is the proton mass. The plasma bubble is optically thick to its thermal photons since initial optical depth is estimated as $\tau_{\rm i} = n_{\rm i}\sigma_{\rm T}R_{\rm i} = 3.7\times10^9(m_{\rm pl}/m_{\rm J})(R_{\rm i}/r_{\rm J})^{-2}$, where $\sigma_{\rm T}$ is the Thomson cross section. Subsequently, the bubble expands due to its thermal pressure. Here, we focus on a simple estimate of the luminosity of the resulting emission by the bubble. For simplicity, we assume that the gas has uniform density and temperature, and a homologous velocity profile.

The bubble is matter dominated, so that the initial temperature is given by

$$T_{\rm i} \sim \frac{GM_*\mu_{\rm p}}{3k_{\rm B}R_*} = 7.7\times10^6\xi_*\ {\rm K} \quad . \tag{17}$$

The initial radiation energy, $E_{\rm rad,i} \sim (aT_{\rm i}^4)(4\pi R_{\rm i}^3/3) \sim 3.8\times10^{43}\xi_*^4(R_{\rm i}/r_{\rm J})^3$erg, where a is the radiation energy constant, is much smaller than $E_{\rm th,i}$ in Eq. (8). The material is initially opaque to its own thermal photons, allowing very little internal energy to escape from its surface. The hot plasma expands adiabatically with an expansion speed comparable to the escape velocity of the star, $v_{\rm esc,*} = (2GM_*/R_*)^{1/2} = 6.2\times10^7\xi_*^{1/2}{\rm cm\ s^{-1}}$. The temperature declines adiabatically as a function of radius R, as $T(R) = T_{\rm i}(R/R_{\rm i})^{-2}$, while the gas number density is given by $n(R) = n_{\rm i}(R/R_{\rm i})^{-3}$.

The expanding bubble becomes optically thin when it cools below the hydrogen recombination temperature of $\sim 10^4$K. The photosphere radius $R_{\rm ph}$ is determined by the condition that the photon diffusion time $t_{\rm diff} = n_{\rm e}\sigma_{\rm T}R^2/c$, where $n_{\rm e} = n_{\rm e}(R)$ is the number density of free electrons, is equal to the expansion time $t_{\rm exp} = R/v_{\rm esc,*}$ at $R_{\rm ph}$. We define $x = n_{\rm e}/n(R)$ as the ionization degree at radius R. The Saha equation for hydrogen,

$$\frac{1-x}{x^2} = n(R)\left(\frac{h^2}{2\pi\mu_{\rm e}kT(R)}\right)^{3/2}\exp\left(\frac{13.6\ {\rm eV}}{kT(R)}\right) \quad , \tag{18}$$

combined with $n(R)/n_{\rm i} = [T(R)/T_{\rm i}]^{3/2}$ and $t_{\rm diff}/t_{\rm exp} = x\tau_{\rm i}(T/T_{\rm i})(v_{\rm esc,*}/c) = 1$, can provide the equation for the temperature $T_{\rm ph}$ at the photosphere, to numerically find $T_{\rm ph} \approx 7200$ K. Due to the exponential term in Eq. (18), the value of $T_{\rm ph}$ is almost independent of the initial state of the bubble. The photosphere radius is therefore,

$$\begin{aligned} R_{\rm ph} &= R_{\rm i}\left(\frac{T_{\rm i}}{T_{\rm ph}}\right)^{1/2} \\ &= 2.3\times10^{11}\xi_*^{1/2}\left(\frac{R_{\rm i}}{r_{\rm J}}\right)\left(\frac{T_{\rm ph}}{7200\ {\rm K}}\right)^{-1/2}\ {\rm cm} \quad , \end{aligned} \tag{19}$$

and the ionization degree at $R = R_{\rm ph}$ is $x(R_{\rm ph}) = 1.4\times10^{-4}(m_{\rm pl}/m_{\rm J})^{-1}(R_{\rm i}/r_{\rm J})^2\xi_*^{-1/2}(T_{\rm ph}/7200\ {\rm K})^{-1}$. The observer would detect blackbody radiation with a temperature $T_{\rm ph}$ and a peak bolometric luminosity,

$$\begin{aligned} L_{\rm p} &= 4\pi R_{\rm ph}^2\sigma T_{\rm ph}^4 \\ &= 1.0\times10^{35}\xi_*\left(\frac{R_{\rm i}}{r_{\rm J}}\right)^2\left(\frac{T_{\rm ph}}{7200\ {\rm K}}\right)^3\ {\rm erg\ s^{-1}} \quad , \end{aligned} \tag{20}$$

Table 1. Predicted optical/infrared peak flux density of the prompt emission. The unabsorbed observed peak flux, $F_\nu^{\rm p}$, is for the distance $d = 10$ kpc from the source with $\xi_* = 1$, $R_{\rm i} = r_{\rm J}$, and $T_{\rm ph} = 7200$ K.

Filter	λ [nm]	ν [10^{14}Hz]	$F_\nu^{\rm p}$ [mJy]
g'	475	6.3	0.97
r'	622	4.8	1.2
i	763	3.9	1.2
y	1020	2.9	1.1
J	1220	2.5	0.91
H	1630	1.8	0.66
K	2190	1.4	0.44
L	3450	0.87	0.21
M	4750	0.63	0.12
N	10500	0.29	0.028

where σ is the Stefan-Boltzmann constant. The peak flux density at frequency $\nu = \nu_{14}\times10^{14}$Hz and a distance $d = 1d_{\rm kpc}$ kpc from the source is then

$$\begin{aligned} F_\nu^{\rm p} &= \frac{L_{\rm p}}{4\pi d^2}\frac{15}{\pi^4\nu}\left(\frac{h\nu}{kT_{\rm ph}}\right)^3 f\left(\frac{h\nu}{kT_{\rm ph}}\right) \\ &= 38\ \xi_*\left(\frac{R_{\rm i}}{r_{\rm J}}\right)^2\nu_{14}^2 d_{\rm kpc}^{-2} f\left(\frac{h\nu}{kT_{\rm ph}}\right)\ {\rm mJy} \quad , \end{aligned} \tag{21}$$

where $f(y) = y(e^y-1)^{-1}$. Table 1 provides the flux density in various observation bands. Note that since $\xi_* \approx 1$ for $M_* \lesssim 2M_\odot$, the observed flux hardly depends on stellar properties. Note that in deriving Eqs. (20), (21) and values in Table 1, we assume that the gas initially expands adiabatically, that is $T(R) \propto [n(R)]^{2/3}$. This relationship is only valid for a monatomic gas of a fixed ionization state. The process of recombination releases energy, which keeps the gas more isothermal than predicted by this relationship. This increases the photosphere radius from that estimated in Eq. (19), increasing the luminosity of the prompt emission.

The typical duration of the transient is comparable to the dynamical time-scale,

$$\Delta T \sim \frac{R_{\rm ph}}{v_{\rm esc,*}} = 3.7\times10^3\left(\frac{R_{\rm i}}{r_{\rm J}}\right)\left(\frac{T_{\rm ph}}{7200\ {\rm K}}\right)^{-1/2}\ {\rm s} \quad , \tag{22}$$

so that the total emission energy is

$$E_{\rm rad} \sim L_{\rm p}\Delta T = 3.7\times10^{38}\xi_*\left(\frac{R_{\rm i}}{r_{\rm J}}\right)^3\left(\frac{T_{\rm ph}}{7200\ {\rm K}}\right)^{5/2}\ {\rm erg} \quad , \tag{23}$$

which is much smaller than the initial internal energy of the bubble. Hence, almost all the initial energy transforms to kinetic energy and gets dissipated when the expanding material interacts with the CSM.

At the moment of tidal disruption, the planet is expected to be vertically collapsed (e.g., Kobayashi et al. 2004; Guillochon et al. 2009). The work done by the tidal force from the star is estimated to be $W \sim (GM_*m_{\rm pl}/R_{\rm p}^2)(r_{\rm pl}/R_{\rm p})(r_{\rm pl}/2) \sim \beta^3Gm_{\rm pl}^2/2r_{\rm pl}$. If the collapsed matter is thermalized and half of this energy is released, then the initial thermal energy becomes $E_{\rm th,i} \sim 8.6\times10^{42}\beta^3(m_{\rm pl}/m_{\rm J})^{5/3}$ erg, where we use an approximate relation, $(r_{\rm pl}/r_{\rm J}) \approx (m_{\rm pl}/m_{\rm J})^{1/3}$ (Rappaport et al. 2013). This is about two orders of magnitude smaller than the energy considered in Eq. (8).

The initial temperature is then estimated as $T_{\rm i} \sim 3.7 \times 10^4 \beta^3 (m_{\rm pl}/m_{\rm J})^{2/3}$K, which is of order the recombination temperature. The bubble expands at a speed comparable to the free-fall velocity of the planet, $v_{\rm ff,pl} \sim (Gm_{\rm pl}/r_{\rm pl})^{1/2} = 4.3 \times 10^6 (m_{\rm pl}/m_{\rm J})^{1/2}(r_{\rm pl}/r_{\rm J})^{-1/2}$cm s^{-1}. Similarly to the previous calculation, we derive a temperature of $T_{\rm ph} \approx 8100$ K at the photosphere radius $R_{\rm ph} = 1.5 \times 10^{10}(m_{\rm pl}/m_{\rm J})^{2/3}\beta^{3/2}(T_{\rm ph}/8100\ {\rm K})^{-1/2}$. The peak bolometric luminosity is then $L_{\rm p} \sim 6.7 \times 10^{32}\beta^3(m_{\rm pl}/m_{\rm J})^{4/3}(T_{\rm ph}/8100\ {\rm K})^3$erg s^{-1}. This is smaller than the main prompt emission (see Eq. 5), however, it could be larger for larger β and/or $m_{\rm pl}$, in which case the emission could be detectable as a precursor arising before the prompt emission.

4 RADIO AFTERGLOW

The expanding plasma maintains a constant velocity $v_{\rm esc,*}$ out to the deceleration radius,

$$R_{\rm dec} = \left(\frac{3m_{\rm pl}}{4\pi\mu_{\rm p}n_{\rm c}}\right)^{1/3} = 6.5{\times}10^{17} n_{\rm c}^{-1/3}\left(\frac{m_{\rm pl}}{m_{\rm J}}\right)^{1/3} {\rm cm}\ , \quad (24)$$

where $n_{\rm c}$ is the density of the CSM. The bubble reaches this radius after $t_{\rm dec} = R_{\rm dec}/v_{\rm esc,*} = 3.3 \times 10^2 \xi_*^{-1} n_{\rm c}^{-1/3}(m_{\rm pl}/m_{\rm J})^{1/3}$yr. During the expansion, the flow interacts with the CSM, generating an external shock with a Mach number of ~ 60. Electron acceleration at the shock results in radio synchrotron emission. The emission lasts until the shock velocity declines to $\sim 1 \times 10^7$ cm s^{-1}, below which the ionization of the acceleration region drops rapidly (Shull & McKee 1979), so that wave damping due to collisions with neutral atoms prevents electrons from being accelerated at the shock front (Drury et al. 1996; Bykov et al. 2000). Assuming the Sedov solution after a time $t_{\rm dec}$ ($R \propto t^{2/5}$ and $v \propto t^{-3/5}$), we estimate the epoch $t_{\rm end}$ at which particle acceleration ceases to be $t_{\rm end} \sim 6.9{\times}10^3 \xi_*^{-1/6} n_{\rm c}^{-1/3}(m_{\rm pl}/m_{\rm J})^{1/3}$yr, corresponding to the radius $R_{\rm end} \sim 2.2 \times 10^{18}\xi_*^{1/3} n_{\rm c}^{-1/3}(m_{\rm pl}/m_{\rm J})^{1/3}$cm. When the acceleration stops, the high-energy electrons starts to escape from the shocked region with an escape time of $\sim$ 1–10 yr, resulting in rapid fading of the emission.

Next we provide a simple estimate of the observed flux and surface brightness of the radio synchrotron emission from the expanding bubble of radius R. The total number of nonthermal electrons, $N_{\rm e}$, is a fraction $\eta_{\rm e}$ of the number of particles originating from the upstream region of the shock over a dynamical time $t_{\rm dyn} = R/v$, where $v \sim v_{\rm esc,*}$ is the expansion speed, that is, $N_{\rm e} = \eta_{\rm e}(4\pi R^2 n_{\rm c} v)t_{\rm dyn} = 4\pi\eta_{\rm e} n_{\rm c} R^3$. We assume a single power-law form of the nonthermal electron distribution, $N(\gamma) \propto \gamma^{-p}$, for $\gamma_{\rm m} < \gamma < \gamma_{\rm M}$, where γ is the electron Lorentz factor. If the dynamical time is sufficiently long ($t_{\rm dyn} \gtrsim 2{\times}10^3$yr), the maximum Lorentz factor $\gamma_{\rm M}$ is determined by the balance of the acceleration time and the synchrotron cooling time (e.g., Yamazaki et al. 2004, 2006, 2015), yielding $\gamma_{\rm M} \approx 1 \times 10^8(\xi_*/fB_{-5})^{1/2}(v/v_{\rm esc,*})$, where $B_{-5} = (B/10\ \mu{\rm G})$ is the post-shock magnetic field strength, and f is a numerical factor of order unity which is determined by the properties of scattering waves and shock geometry at the acceleration site. Thus, we find that radio-emitting electrons have a much smaller Lorentz factor than $\gamma_{\rm M}$.

Table 2. Radio surface brightness at a radius $R = 10^{18}$cm and frequency $\nu = 1$ GHz for power-law electron distribution with index p. Other parameters are taken as $n_{\rm c} = \eta_{-5} = B_{-5} = \gamma_{\rm m} = 1$.

p	$(\nu/\nu_{\rm m})^{(1-p)/2}$	$S_{\nu=1{\rm GHz}}$ [Jy sr^{-1}]
2.0	1.7×10^{-4}	2.0×10^5
2.1	7.0×10^{-5}	8.4×10^4
2.2	2.9×10^{-5}	3.5×10^4
2.3	1.2×10^{-5}	1.5×10^4
2.4	5.2×10^{-6}	6.2×10^3
2.5	2.2×10^{-6}	2.6×10^3
3.0	2.8×10^{-8}	34

The synchrotron cooling time of radio emitting electrons is much longer than the dynamical time. Hence, the observed flux density at frequency ν is (e.g., Sari, Piran, & Narayan 1998)

$$F_\nu \sim \frac{N_{\rm e}}{4\pi d^2}\frac{\mu_{\rm e}c^2\sigma_{\rm T}B}{3e}\left(\frac{\nu}{\nu_{\rm m}}\right)^{(1-p)/2} = 4.0 \times 10^2 \frac{n_{\rm c}\eta_{-5}B_{-5}}{d_{\rm kpc}^2}\left(\frac{R}{10^{18}{\rm cm}}\right)^3\left(\frac{\nu}{\nu_{\rm m}}\right)^{(1-p)/2} {\rm Jy}\ , \quad (25)$$

where $\eta_{-5} = (\eta_{\rm e}/10^{-5})$ and $\nu_{\rm m} = 28B_{-5}\gamma_{\rm m}^2$Hz is the characteristic synchrotron frequency from electrons with a minimum Lorentz factor $\gamma_{\rm m}$. Note that for our parameters, $(\nu/\nu_{\rm m})^{(1-p)/2}$ is much lower than unity (see Table 2). The surface brightness, S_ν, is the flux density divided by the solid angle of the source, $\Omega \sim \pi(R/d)^2$, yielding

$$S_\nu \sim 1.2 \times 10^9 n_{\rm c}\eta_{-5}B_{-5}\frac{R}{10^{18}{\rm cm}}\left(\frac{\nu}{\nu_{\rm m}}\right)^{(1-p)/2} {\rm Jy\ sr}^{-1}\ . \quad (26)$$

Table 2 shows the results for different values of p with fixed parameters, $n_{\rm c} = \eta_{-5} = B_{-5} = \gamma_{\rm m} = 1$. For comparison, the surface brightness of the faintest Galactic supernova remnants is $\sim 10^4$Jy sr^{-1} (Arbutina & Urošević 2005), of the same order as the typical value of the diffuse Galactic radio emission (e.g., de Oliveira-Costa et al. 2008).

5 DISCUSSION

5.1 Event rate

The event rate of the transients considered in this paper is highly uncertain. It depends on various processes such as the formation and early dynamics of giant planets. The rate per galaxy is roughly given by $\sim {\rm SFR} \times (\alpha N_{\rm pl})/\langle m\rangle \approx 5(\alpha N_{\rm pl})\ {\rm yr}^{-1}$, where SFR is the star formation rate (see § 1), $N_{\rm pl}$ is the average number of giant planets per star, and α is the fraction of planets that are close enough to directly hit the stellar surface. At present, it is highly uncertain how many giant planets are formed in a circumstellar disk because of the unknown disk mass. Recent numerical simulations have indicated that disks are likely to be very massive so as to harbor a lot of giant planets, suggesting $N_{\rm pl}$ could be ~ 10 (e.g., Machida et al. 2011). The typical value of α is also uncertain. To directly impact the stellar surface, planets must have a penetration factor larger than unity (see § 2). It is possible that planet-planet interactions increase the eccentricity making the pericenter small enough for a direct hit on the star (Li et al. 2014). However, tidal dissipation would tend to limit the eccentricity growth when the pericenter becomes small. Another regime allowing for a direct collision between the planet and the star emerges in the early phase of the circumstellar disk, when the cloud core is still collapsing and the motion of the newly formed planet is chaotic. Some planets are formed at a large radii ($\sim$ 20–50 AU) and orbit only for 2–3 orbital periods, after which they fall into the star (Machida et al. 2011). In this case, the eccentricity of the final orbit before the impact on the star may be near unity. Here we expect that $\alpha N_{\rm pl}$ would be of order unity or even larger. Future observation of the rate of the transients will constrain these highly uncertain scenarios and shed light on the process of the planet formation. In particular, if the event rate is high, we will be able to confirm that the circumstellar disk is massive.

It is also possible that after ingestion of Jovian planets, stars become metal rich (e.g., Sandquist et al. 1998, 2002; Cody & Sasselov 2005). One can get an upper limit on the event rate from the fraction of metal rich stars. However, the high metallicities could result from other processes (e.g. inhomogeneous supernova enrichment of the interstellar medium). The actual rate can be calibrated as a fraction of this maximum rate.

5.2 Prospects for future observations

While the prompt emission flare from the tidal disruption of a planet cannot be detected in the optical band due to Galactic dust extinction, the unabsorbed flux in the K-band is ~ 0.2 mJy at the distance of 10 kpc, potentially detectable with infrared sky surveys, such as UKIRT Infrared Deep Sky Survey (UKIDSS: Hewett et al. 2006; Lawrence et al. 2007) and VISTA Variables in the Via Lactea (VVV: Minniti et al. 2010). If $\alpha N_{\rm pl} \lesssim 1$, the expected event rate $\sim 5(\alpha N_{\rm pl})$ events yr^{-1} for the entire Galactic plane (see § 1), is too small for detectability by current transient surveys. The duration ΔT of the expected prompt flares could be comparable to that of superflares of stars (e.g., Schaefer et al. 2000; Maehara et al. 2012). The total emission energy in the optical band is about an order of magnitude larger than the largest superflares (Shibayama et al. 2013; Balona 2015).

The Andromeda galaxy, M31, located at the distance of 0.78 Mpc (Stanek & Garnavich 1998), has a similar mass and star formation rate to the Milky Way (Williams 2003), and hence a similar event rate of the transients proposed in this paper. Unlike the Milky Way case, telescopes with a large field of view can cover the entire volume of M31. The expected AB magnitudes are 25.9 and 25.7 mag at g and r bands, respectively, for our fiducial parameters. The fluxes become higher for planets more larger than Jupiter. Using the pixel lensing technique (i.e., differencial image photometry: Crotts 1992; Baillon et al. 1993; Tomaney & Crotts 1996; Calchi Novati 2010), those transients could be detected by future instruments with better sensitivity like Subaru Hyper Suprime-Cam[3].

[3] http://www.naoj.org/Projects/HSC/index.html

Events in the Large Magellanic Cloud (LMC) could also be detected. For the distance of 48.5 kpc to LMC (Macri et al. 2006), the expected AB magnitudes are 19.9 and 19.6 mag at g and r bands respectively for our fiducial parameters, which are detectable by current surveys like PAndromeda (Lee et al. 2012). The event rate is only slightly smaller than M31 because the total star formation rate of the LMC is $\approx 0.4 M_\odot$ yr^{-1} (Harris & Zaritsky 2009).

If the electron index p is smaller than about 2.3, the radio afterglow would be detectable for $\sim 10^{3-4}$yr after the disruption event. The radio surface brightness is expected to be lower than young supernova remnants because the magnetic field is weak. As a result, the surface brightness and diameter of the source would be distinguishable from the values expected for supernova remnants (Arbutina & Urošević 2005). Note that the classical nova eruption ejects typically 10^{-5} to $10^{-4} M_\odot$ of matter (e.g., Bode 2010; Roy et al. 2012), which is one or two orders of magnitude smaller than that assumed in our present model, $m_{\rm pl} \sim 10^{-3} M_\odot$. The observed radio emission of novae is typically well characterized by the free-free emission process, which peaks at ~ 1 yr after the eruption (e.g., Ribeiro et al. 2014). After that, the observed radio emission declines, and no radio synchrotron halo has been ever detected (e.g., Ribeiro et al. 2014). Since the ejecta mass of novae and our planet-star collision model are close, the current radio upper limits on the synchrotron shock emission may potentially constrain the present model. Equations (25) and (26) imply that the flux F_ν and the surface brightness S_ν increase until the deceleration time $t_{\rm dec}$, and takes maximum at $R = R_{\rm dec}$ if other model parameters are constant with time. Since $R_{\rm dec} \propto M_{\rm ej}^{1/3}$ where $M_{\rm ej}$ is the ejecta mass (that is, $M_{\rm ej} \approx m_{\rm pl} \sim 10^{-3} M_\odot$ for our present model, and $M_{\rm ej} \sim 10^{-5} - 10^{-4} M_\odot$ for the nova eruption), we get $F_\nu \propto R_{\rm dec}^3 \propto M_{\rm ej}$ and $S_\nu \propto R_{\rm dec} \propto M_{\rm ej}^{1/3}$, so that the radio synchrotron afterglows of the classical novae show surface brightness several times smaller than those of the star-planet collision. As shown in Table 2, our present synchrotron shock model predicts low surface brightness that is comparable to the present-day detection limit. This may explain why the related radio synchrotron afterglow has not been detected as of yet.

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